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F-term uplifting and moduli stabilization consistent with Kähler invariance

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Abstract: An important ingredient in the construction of phenomenologically viable superstring models is the uplifting of Anti-de Sitter supersymmetric critical points in the moduli sector to metastable Minkowski or de Sitter vacua with broken supersymmetry. In all cases described so far, uplifting results in a displacement of the potential minimum away from the critical point and, if the uplifting is large, can lead to the disappearance of the minimum altogether. We propose a variant of F-term uplifting which exactly preserves supersymmetric critical points and shift symmetries at tree level. In spite of a direct coupling, the moduli do not contribute to supersymmetry breaking. We analyse the stability of the critical points in a toy one-modulus sector before and after uplifting, and find a simple stability condition depending solely on the amount of uplifting and not on the details of the uplifting sector. There is a region of parameter space, corresponding to the uplifting of local AdS *maxima* –or, more importantly, local minima of the Kähler function– where the critical points are stable for *any* amount of uplifting. On the other hand, uplifting to (non-supersymmetric) Minkowski space is special in that all SUSY critical points, that is, for *all* possible compactifications, become stable or neutrally stable.

Keywords: Supergravity Models, Flux compactification, dS vacua in string theory.

Contents

1. Introduction

The uplifting of supersymmetric critical points from Anti-de Sitter to Minkowski or de Sitter vacua is a crucial but still not completely understood element in the standard KKLT scenario of moduli stabilization in type IIB string theory [1]. The dilaton and complex structure moduli are stabilized by fluxes while other non-perturbative effects stabilize the remaining Kähler moduli at constant values that preserve supersymmetry – leading to a cosmological AdS vacuum with negative potential energy. In the original model, uplifting to a positive value of the potential is achieved by anti-D3 branes which break supersymmetry explicitly. Subsequent work has concentrated mainly on D-term [2, 3, 4, 5, 6, 7, 8, 9, 10, 11] and F-term uplifting [12, 13, 14, 15, 16, 17, 18, 19, 20, 21, 22, 23] as interesting alternatives where there is an explicit supergravity description and supersymmetry is only broken spontaneously, which gives better calculational control. Uplifting by Kähler corrections has also been considered [24, 25].

In the case of F-term uplifting, one possible strategy is to combine the moduli with another sector whose SUSY breaking properties and phenomenology are known in isolation (Polonyi model [26], O'Raifertaigh [27], ISS [28]) and hope -or, rather, check- that the interaction with the moduli will respect the basic features of both sectors. There are examples in the literature both with and without direct couplings in the superpotential between the modulus and the SUSY breaking sector.

Intuitively, one would expect compactification to be a high energy phenomenon, possibly near the Planck scale, and therefore decoupled from the low energy effective action

describing our current Universe. This is certainly our experience. In spite of a plethora of very precise cosmological and accelerator data, we still see no evidence of extra dimensions. Upcoming experiments such as the Large Hadron Collider at CERN or the Planck mission may change this picture but in any case the effect is so small that it still makes sense to look for a general framework in which at least some of the moduli are completely stabilized and as decoupled as possible from phenomena far below the compactification scale. Since gravity couples to all fields and supersymmetry restricts the form of the interactions, this task has proved somewhat tricky in supergravity.

The broader question we revisit here is how to couple two supergravity sectors in such a way that they interact *as little as possible*. We must stress that this is not a well-defined condition, as the answer depends strongly on what properties we wish to preserve. From the point of view of low-energy phenomenology, requiring gravitational-strength couplings may be sufficient; however, at higher energies this condition can become difficult to check explicitly when there are moduli or inflatons involved with near-Planckian vacuum expectation values. Instead, in these regimes, supersymmetry seems a much more powerful guiding principle and one that has been very successful in other contexts. It also facilitates comparison with string theory, where the supersymmetry of some configurations can be determined explicitly without reference to N=1 supergravity.

A particularly challenging decoupling problem is encountered when trying to construct stringy models of slow roll inflation (for a recent review see [29] and references therein). In general it is impossible to know if a given field is a good candidate for an inflaton until the complete potential is known because the fields always evolve in the steepest direction of the potential. Even if the slow-roll conditions are satisfied for a given field one must make sure that all other fields are properly stabilized. To make matters worse, in the standard KKLT and racetrack [30, 31, 32] scenarios, inflation can easily destabilize the moduli, leading to decompactification. It is therefore important to understand what kind of supergravity lagrangians have interactions between the inflaton and the moduli such that, on the one hand, the slow-roll conditions are not spoilt and, on the other, no modulus becomes unstable during inflation. Shift symmetries are sometimes invoked to solve the first problem, for instance in relation to BPS configurations of D3-D7 branes [33, 34, 35][1]. The second condition, in the absence of finetuning, seems to lead to the generic constraint that the scale of inflation has to be below the gravitino mass [37, 38]. Therefore it is important to find phenomenologically viable models with this property or else find ways of evading the constraint.

In this paper we propose a way of coupling supergravity sectors that preserves some of their supersymmetry properties such as supersymmetric critical points and shift symmetries, at tree level. A very important clue comes from Kähler invariance because the properties we wish to preserve are invariant under Kähler transformations. We will require the total action of the coupled sectors to be invariant under Kähler transformations of the individual sectors. This limits the validity of our approach since each sector must have an independent description in terms of N=1 SUGRA, with a non-zero superpotential, and

[1]Shift symmetries have also been used to solve the η problem in more general contexts [36].

we assume all fields are Kähler invariant. The assumption of non-zero superpotentials is reasonable since our world, the visible sector, is described by a near–Minkowski vacuum with broken supersymmetry, and it also holds generically for the moduli sector [1]. Moreover, $D-$term uplifting is not possible without $F-$term uplifting as well [39, 40], so we concentrate here on the simplest case of uncharged chiral superfields, say $\{z^\alpha\}$ for the supersymmetric ("moduli") sector and $\{\phi^i\}$ for the SUSY-breaking sector responsible for the uplifting. In that case the D-terms are zero and the two sectors are fully described by Kähler functions[2] $G^{(1)}(z,\bar{z}) = K^{(1)}(z,\bar{z}) + \ln|W^{(1)}(z)|^2$ and $G^{(2)}(\phi,\bar{\phi}) = K^{(2)}(\phi,\bar{\phi}) + \ln|W^{(2)}(\phi)|^2$. $K^{(1)}$ and $K^{(2)}$ are, as usual, the Kähler potentials that determine the scalar manifold metric and $W^{(1)}$ and $W^{(2)}$ the holomorphic superpotentials. The condition of Kähler invariance then tells us that, if the Kähler potential K of the coupled system is of the form

$$K = f\ (\ K^{(1)}, K^{(2)}\ ) \qquad \text{or, more generally,} \qquad F\ (\ K, K^{(1)}, K^{(2)}) = 0 \tag{1.1}$$

for some function F, the Kähler function of the coupled system G must be of the form

$$G = f\ (\ G^{(1)}, G^{(2)}\ ) \qquad \text{or, more generally,} \qquad F\ (\ G, G^{(1)}, G^{(2)}\ ) = 0\ \ . \tag{1.2}$$

In the particular case where the kinetic terms of the two sectors are decoupled and the Kähler potential is separable,

$$K(z^\alpha, \bar{z}^{\bar{\alpha}}, \phi^i, \bar{\phi}^{\bar{i}}) = K^{(1)}(z^\alpha, \bar{z}^{\bar{\alpha}}) + K^{(2)}(\phi^i, \bar{\phi}^{\bar{i}})\ \ ,$$

this prescription leads uniquely to the ansatz

$$G(z^\alpha, \bar{z}^{\bar{\alpha}}, \phi^i, \bar{\phi}^{\bar{i}}) = G^{(1)}(z^\alpha, \bar{z}^{\bar{\alpha}}) + G^{(2)}(\phi^i, \bar{\phi}^{\bar{i}})\ , \tag{1.3}$$

that is, to the *product* (as opposed to the *sum*) of superpotentials

$$K(z^\alpha, \bar{z}^{\bar{\alpha}}, \phi^i, \bar{\phi}^{\bar{i}}) = K^{(1)}(z^\alpha, \bar{z}^{\bar{\alpha}}) + K^{(2)}(\phi^i, \bar{\phi}^{\bar{i}}) \tag{1.4}$$

$$W(z^\alpha, \phi^i) = W^{(1)}(z^\alpha) W^{(2)}(\phi^i). \tag{1.5}$$

This ansatz is not new. Binetruy *et al.* [41] discuss it as a sufficient condition for integrating out heavy chiral multiplets in a supersymmetric way. Later Hsu *et al.* [33] used this ansatz to characterize an effective SUGRA theory describing D3-D7 brane inflation in a type IIB string compactification.

By contrast, the usual ansatz invoked for gravitational strength couplings,

$$K(z^\alpha, \bar{z}^{\bar{\alpha}}, \phi^i, \bar{\phi}^{\bar{i}}) = K^{(1)}(z^\alpha, \bar{z}^{\bar{\alpha}}) + K^{(2)}(\phi^i, \bar{\phi}^{\bar{i}}) \tag{1.6}$$

$$W(z^\alpha, \phi^i) = W^{(1)}(z^\alpha) + W^{(2)}(\phi^i). \tag{1.7}$$

suffers from an ambiguity in the case where the superpotentials are nonzero, since it depends on the Kähler gauge chosen in each sector before combining them. A Kähler transformation

[2] We use units with $M_P = 1$ throughout

of each sector separately, $K^{(I)} \to K^{(I)} + 2\,\mathrm{Re} f^{(I)}$, and $W^{(I)} \to W^{(I)} e^{-f^{(I)}}$, $I = 1, 2$, leads to

$$K = K^{(1)}(z,\bar{z}) + K^{(2)}(\phi,\bar{\phi}) \to K + 2\mathrm{Re}(f^{(1)}(z) + f^{(2)}(\phi)) \tag{1.8}$$

$$W = W^{(1)}(z) + W^{(2)}(\phi) \to W^{(1)}(z)e^{-f^{(1)}(z)} + W^{(2)}(\phi)e^{-f^{(2)}(\phi)} \tag{1.9}$$

which is equivalent to

$$K = K^{(1)}(z,\bar{z}) + K^{(2)}(\phi,\bar{\phi}) \tag{1.10}$$

$$W = W^{(1)}(z)e^{f^{(2)}(\phi)} + W^{(2)}(\phi)e^{f^{(1)}(z)}\,, \tag{1.11}$$

a completely different final theory with direct couplings between the two sectors. The relation between the ansatz (1.7) and gravitational strength couplings is therefore more subtle than is usually assumed.

As we mentioned before, the ansatz that we propose to couple sectors (1.3) exactly preserves supersymmetric critical points, in contrast with the standard ansatz (1.7), which generically leads to a shift of these points. If we take z_0^α and ϕ_0^i to be supersymmetric critical points of the z and ϕ sectors respectively:

$$[\partial_\alpha W^{(1)} + \partial_\alpha K^{(1)} W^{(1)}]_{z_0^\alpha} = 0 \qquad [\partial_i W^{(2)} + \partial_i K^{(2)} W^{(2)}]_{\phi_0^i} = 0, \tag{1.12}$$

the field configuration (z_0^α, ϕ_0^i) in general will not be a SUSY critical point of the total theory defined by (1.7):

$$[\partial_\alpha W + \partial_\alpha K W]_{z_0^\alpha,\phi_0^i} = \partial_\alpha K^{(1)} W^{(2)}|_{z_0^\alpha,\phi_0^i} \qquad [\partial_i W + \partial_i K W]_{z_0^\alpha,\phi_0^i} = \partial_i K^{(2)} W^{(1)}|_{z_0^\alpha,\phi_0^i} \tag{1.13}$$

In order to preserve the supersymmetric critical points additional conditions must be imposed, either the superpotentials of the individual sectors vanish at the critical point $W^{(1)}|_{z_0^\alpha} = W^{(2)}|_{\phi_0^i} = 0$ or the first derivatives of the Kähler potential $\partial_\alpha K^{(1)}|_{z_0^\alpha} = \partial_i K^{(2)}|_{\phi_0^i} = 0$ are zero at the critical point, or some other suitable combination that makes both F-terms zero. The moduli sectors appearing in the KKLT framework generically lead to a SUSY critical point where the superpotential does not vanish, but the second condition can be satisfied provided an explicit Kähler transformation is performed before the superpotentials are added.

The paper is organized as follows. In section 2 we will introduce our notation while reviewing some basic features of $\mathcal{N} = 1$ SUGRA actions. In section 3 we study the coupling of two sectors following the ansatz (1.3) and its basic, model-independent properties. We are interested in applications to F-term uplifting so we consider a supersymmetric sector described by an arbitrary Kähler function admitting one or more critical points, which are also critical points of the potential. The uplifting sector is also arbitrary except for the requirement that it must have a SUSY breaking, Minkowski or de Sitter vacuum or plateau at tree level -the latter, e.g. from a shift symmetry-. In section 4 we look at the stability of the uplifted moduli in the simplest possible case: a toy model consisting of one supersymmetric “modulus” field coupled to a supersymmetry breaking “uplifting” sector consisting of neutral scalar fields. We conclude with a summary of the main results in section 5.

2. Review of $\mathcal{N}=1$ supergravity

We start with a quick review of the relevant SUGRA formulae to fix our notation. We take $M_{\text{Planck}}=1$. Consider an $\mathcal{N}=1$ SUGRA sector that we will call the *supersymmetric sector* consisting of neutral chiral superfields $\{z^\alpha\}$. It is described by a Kähler potential $K(z,\bar z)$ and a superpotential $W(z)$. The kinetic terms are

$$\int d^4x\sqrt{-g}\; K_{\alpha\bar\beta}\;\partial_\mu z^\alpha\partial_\nu \bar z^{\bar\beta} g^{\mu\nu}.$$

We will use the standard notation denoting derivatives by subscripts:

$$\partial_\alpha K\equiv K_\alpha \qquad \partial_{\bar\beta}K\equiv K_{\bar\beta} \quad \partial_{\alpha\bar\beta}K\equiv K_{\alpha\bar\beta} \quad \text{etc}\ldots, \tag{2.1}$$

and the indices being raised and lowered with the Kähler metric $K_{\alpha\bar\beta}$ and its inverse $K^{\alpha\bar\beta}=K^{-1}_{\alpha\bar\beta}$. Since the fields are uncharged and there are no gauge fields, the D-terms are zero and the potential is given by

$$V=e^K[K^{\alpha\bar\beta}(\partial_\alpha W+\partial_\alpha K\,W)(\partial_{\bar\beta}\bar W+\partial_\beta K\,\bar W)-3|W|^2]\ . \tag{2.2}$$

In what follows we shall omit the superscripts α and i of the fields. The action and the supersymmetry transformations are invariant under Kähler transformations,

$$K(z,\bar z)\;\rightarrow\;K(z,\bar z)+f(z)+\bar f(\bar z) \tag{2.3}$$

$$W(z)\;\rightarrow\;W(z)e^{-f(z)} \tag{2.4}$$

where $f(z)$ is an arbitrary holomorphic function. If $W\neq 0$, both can be expressed in terms of the *Kähler function*,

$$G(z,\bar z)=K(z,\bar z)+\ln|W(z)|^2, \tag{2.5}$$

which is invariant under Kähler transformations. In particular, since $G_{\alpha\bar\beta}=K_{\alpha\bar\beta}$, the kinetic term T and potential V can be written as:

$$T=G_{\alpha\bar\beta}\partial_\mu z^\alpha\partial_\nu\bar z^{\bar\beta}g^{\mu\nu} \qquad V=e^G[G^{\alpha\bar\beta}G_\alpha G_{\bar\beta}-3] \tag{2.6}$$

The points, $z=z_0$, where the F-terms vanish,

$$D_\alpha W|_{z=z_0}=\partial_\alpha W|_{z=z_0}+\partial_\alpha K|_{z=z_0}W|_{z=z_0}=0 \qquad\Leftrightarrow\qquad \partial_\alpha G|_{z=z_0}=0$$

are called SUSY critical points. They are automatically critical points of V because

$$\partial_\gamma V|_{z=z_0}=\left[G_\gamma V+e^G\partial_\gamma(G^{\alpha\bar\beta}G_\alpha G_{\bar\beta})\right]|_{z=z_0}=0 \tag{2.7}$$

Unlike in global SUSY, where supersymmetric critical points are always absolute minima of V, the critical points in SUGRA may be local minima, maxima or saddle points. In SUGRA, supersymmetric critical points are always[3] AdS since $V(z_0)=-3e^{G(z_0)}$. This

[3]The case $e^{G(z_0)}=0$, which corresponds with a Minkowski spacetime is excluded by the condition $W\neq 0$

means that local maxima or saddle points are not necessarily unstable before uplifting [42]. However, after uplifting to Minkowski or dS, only local minima are stable. In a Minkowski background, the gravitino mass is $m_{3/2}^2 = e^G$.

In the next section we consider uplifting to positive V by coupling the $\{z^\alpha\}$ fields to another set of fields $\{\phi^i\}$ that we shall name the *uplifting* sector, also composed of neutral chiral superfields. Ultimately, the visible sector must also be included but this is beyond the scope of this paper. Here we are interested in the effect of uplifting on the moduli.

3. F-term uplifting consistent with Kähler invariance

We consider the coupling of two sectors with neutral chiral superfields $\xi^I = z^\alpha, \phi^i$. We assume that each sector has a SUGRA description with a well-defined Kähler function (a non-zero superpotential). If the sectors are sufficiently decoupled we expect the kinetic terms to add at tree level without interaction, so we take

$$K = K^{(1)}(z,\bar{z}) + K^{(2)}(\phi,\bar{\phi}) \tag{3.1}$$

which makes the Kähler metric block diagonal, and thus the kinetic terms decouple:

$$K_{I\bar{J}}\partial_\mu \xi^I \partial^\mu \bar{\xi}^{\bar{J}} = K^{(1)}_{\alpha\bar{\beta}}(z,\bar{z})\partial_\mu z^\alpha \partial^\mu \bar{z}^{\bar{\beta}} + K^{(2)}_{i\bar{j}}(\phi,\bar{\phi})\partial_\mu \phi^i \partial^\mu \bar{\phi}^{\bar{j}} \tag{3.2}$$

As explained in the introduction, the ansatz (3.1) plus the condition of invariance under Kähler transformations of the individual sectors:

$$\begin{aligned} K^{(1)}(z,\bar{z}) &\to K^{(1)}(z,\bar{z}) + f^{(1)}(z) + \bar{f}^{(1)}(\bar{z}) \qquad & W^{(1)}(z) &\to W^{(1)}(z)e^{-f^{(1)}(z)} \\ K^{(2)}(\phi,\bar{\phi}) &\to K^{(2)}(\phi,\bar{\phi}) + f^{(2)}(\phi) + \bar{f}^{(2)}(\bar{\phi}) \qquad & W^{(2)}(\phi) &\to W^{(2)}(\phi)e^{-f^{(2)}(\phi)} \end{aligned} \tag{3.3}$$

for arbitrary $f^{(1)}(z)$ and $f^{(2)}(\phi)$, forces us to add the full Kähler functions:

$$G(z,\bar{z},\phi,\bar{\phi}) \equiv K + \ln|W|^2 = A(z,\bar{z}) + B(\phi,\bar{\phi}) \tag{3.4}$$

where A and B are the corresponding Kähler functions for both sectors. In our previous notation, $A \equiv G^{(1)} \equiv K^{(1)} + \ln|W^{(1)}|^2$, $B \equiv G^{(2)} = K^{(2)} + \ln|W^{(2)}|^2$. The potential becomes

$$V = e^G[G^{I\bar{J}}G_I G_{\bar{J}} - 3] = e^{A+B}[A^{\alpha\bar{\beta}}A_\alpha A_{\bar{\beta}} + B^{i\bar{j}}B_i B_{\bar{j}} - 3] \tag{3.5}$$

Note that the first term in the square bracket is a function of $(z,\bar{z})$ only, and the uplifting is provided by the second term, $B^{i\bar{j}}B_i B_{\bar{j}}$, which is a function of $(\phi,\bar{\phi})$ alone. The exponential outside the bracket provides a direct coupling between the two sectors. Alternatively, we can write

$$V = e^B V_A(z) + e^A V_B(\phi) + 3e^{A+B} \tag{3.6}$$

where $V_A(z) = e^A[A^{\alpha\bar{\beta}}A_\alpha A_{\bar{\beta}} - 3]$ would be the potential calculated for the z sector alone and similarly for $V_B(\phi)$.

3.1 Critical points and stability; SUSY breaking

As we pointed out in the introduction, coupling the uplifting sector to the supersymmetric sector according to the ansatz (3.4) respects the SUSY properties of the individual sectors. In particular the supersymmetric critical points of the z-sector are still critical points of the full potential. To see this suppose $z = z_0$ is a SUSY critical of the z-sector, $\partial_\alpha A(z_0) = 0$, then from (3.6) we can see that z_0 also satisfies the necessary condition to be a critical point of the full potential:

$$V_\alpha(z_0) = [e^B V_{A\ \alpha} + A_\alpha\, e^A V_B + 3A_\alpha\, e^{A+B}]_{z=z_0} = 0. \tag{3.7}$$

and furthermore the F-terms for z vanish in the full model:

$$|F_z|^2 = e^G G^{\alpha\bar{\beta}} G_\alpha G_{\bar{\beta}}|_{z_0} = 0, \quad \text{since we have} \quad G_\alpha(z_0) = A_\alpha(z_0) = 0, \tag{3.8}$$

which means that the moduli sector does not contribute to SUSY breaking at tree level.

For $z = z_0$ to really correspond to a critical point of the full potential we have to find a configuration $\phi = \phi_0$ so that the criticality condition for the uplifting sector $\partial_i V(z_0, \phi_0) = 0$ is also satisfied. Using that $V_A|_{z=z_0} = -3e^{A(z_0)}$ we find that the full potential evaluated at the point $z = z_0$ is given by the expression

$$V|_{z=z_0} = e^{A(z_0)} V_B(\phi)\ , \tag{3.9}$$

which differs only from the original potential of the uplifting sector by an overall factor $e^{A(z_0)}$. Therefore in order to be at an extremum (local minimum) of the full potential we just have to fix the uplifting sector at any extremum (local minimum) of V_B, which we denote by $\phi = \phi_0$.

Now we turn to the issue of stability of the critical point (z_0, ϕ_0) with z_0 being a supersymmetric critical point of the $z-$sector. An interesting feature of this model is that it is enough to analyze the stability of the critical point along the ϕ^i and z^α directions separately. Indeed, the the mass matrix has a block diagonal form, i.e. $V_{\alpha i}(z_0, \phi_0) = V_{\alpha\bar{i}}(z_0, \phi_0) = 0$, and therefore we just have to check whether the eigenvalues of the matrices

$$\begin{pmatrix} V_{\alpha\bar{\beta}}(z_0,\phi_0) & V_{\alpha\beta}(z_0,\phi_0) \\ V_{\bar{\alpha}\bar{\beta}}(z_0,\phi_0) & V_{\bar{\alpha}\beta}(z_0,\phi_0) \end{pmatrix} \quad \text{and} \quad \begin{pmatrix} V_{i\bar{j}}(z_0,\phi_0) & V_{ij}(z_0,\phi_0) \\ V_{\bar{i}\bar{j}}(z_0,\phi_0) & V_{\bar{i}j}(z_0,\phi_0) \end{pmatrix} \tag{3.10}$$

are all positive. The cross terms, $V_{\alpha i}$ and $V_{\alpha\bar{i}}$, can be calculated taking the derivatives of (3.7) w.r.t. ϕ^i and $\bar{\phi}^{\bar{i}}$. Then, using that A_z and $V_{A\ z}$ are zero at $z = z_0$, it is easy to check that they will all vanish once evaluated at the critical point:

$$\begin{aligned} V_{\alpha\, i}|_{z=z_0} &= [B_i e^B V_{A\ \alpha} + A_\alpha e^A V_{B\ i} + 3e^{A+B} A_\alpha B_i]_{z=z_0} = 0\ , \\ V_{\alpha\, \bar{i}}|_{z=z_0} &= [B_{\bar{i}} e^B V_{A\ \alpha} + A_\alpha e^A V_{B\ \bar{i}} + 3e^{A+B} A_\alpha B_{\bar{i}}]_{z=z_0} = 0. \end{aligned} \tag{3.11}$$

Before we continue discussing the stability of the critical point let us make some remarks about the uplifting of the SUSY critical point of the $z-$sector, z_0. For later convenience let us introduce the notation

$$b(\phi) = B^{i\bar{j}} B_i B_{\bar{j}}, \tag{3.12}$$

so that the potential of the uplifting sector alone, V_B, reads

$$V_B(\phi) = e^{B(\phi)}[b(\phi) - 3]. \tag{3.13}$$

The quantity b is related to the F-terms calculated from the ϕ sector alone $|F_\phi|^2 = e^B\, b$, and therefore to the SUSY breaking scale M_s, since $|F_\phi| = M_s^2$.
In view of equation (3.9), which gives the vacuum expectation value of the full potential with z fixed at z_0, it is now clear that in order to uplift the SUSY critical point to Minkowski or de Sitter, we have to stabilize the $\phi-$sector at Minkowski of de Sitter vacuum of V_B, ϕ_0, so that $V_B|_{\phi=\phi_0} = 0$ or $V_B|_{\phi=\phi_0} > 0$ respectively. Thus for uplifting to Minkowski we need the $\phi-$sector to be stabilized at a point ϕ_0 with $b(\phi_0) = 3$, while for uplifting to de Sitter we have to require $b(\phi_0) > 3$.

As we argued above in order to analyze the stability of the critical point (z_0, ϕ_0) it is enough to study the stability along the z^α and ϕ^i directions separately, since the cross terms of the mass matrix vanish (3.11). Therefore in order to analyze the stability along the z^α directions it is enough to study the potential evaluated at $\phi = \phi_0$, which reads:

$$V|_{\phi=\phi_0} = e^{B(\phi_0)}[V_A(z) + e^{A(z)} b(\phi_0)]. \tag{3.14}$$

From this equation it is clear that the result of the stability analysis will depend on the uplifting sector only through the value of $b(\phi_0)$. A remarkable property of our model is that *all* SUSY critical points of the $z-$sector are either stable or marginally stable in the z^α directions after the uplifting to Minkowski vacuum ($b = 3$). To see this we set $b(\phi_0) = 3$ in the previous equation (3.14), then the full potential evaluated at the point ϕ_0 reads

$$V|_{\phi=\phi_0} = e^{A(z)+B(\phi_0)} A^{\alpha\bar\beta} A_\alpha A_{\bar\beta} \geq 0 \quad \text{for all } z. \tag{3.15}$$

Since, by assumption, $V(z_0, \phi_0) = 0$, the condition (3.15) implies that no fluctuation of the fields on the $z-$sector can decrease the energy, and therefore the point (z_0, ϕ_0) is either a local minimum or a plateau along the z^α directions. A similar result was found in [43]. Here Blanco-Pillado *et al.* argued that SUSY vacua with vanishing cosmological constant are automatically stable, up to flat directions. Note that such minima necessarily have a vanishing superpotential, while in our case we are assuming that the superpotential does not vanish at the critical point. The main difference is that the Minkowski critical point in our model is *not supersymmetric*, but the coupling to the SUSY breaking sector using (1.3) respects the supersymmetric character of the z-sector enough to ensure the stability of the critical point along the z^α directions.

In general we cannot make a similar statement for uplifting to de Sitter critical points, and the stability will depend on the masses of the z^α fields before the uplifting and the value of b. However the analysis of the stability simplifies for large values of the uplifting parameter b. With the total potential written as in (3.14) we can see that for high values of $b(\phi_0)$ the second term dominates, and therefore the minima of e^A are the ones that will survive the uplifting. Moreover, the higher the value of b the higher the masses of the z^α

fields will be after the uplifting. Note also that minima of e^A are not necessarily minima of V_A. In the one-modulus example described in section 4 the minima of e^A correspond to either local maxima or saddle points of the potential V_A.

We will now comment on the stability of the critical point (z_0, ϕ_0) along the ϕ^i directions. Since the stability analysis along the ϕ^i directions is decoupled from the one along the z^α directions, it is enough to consider the potential once evaluated at $z = z_0$ (3.9). In view of this equation we can conclude that the minima of the combined potential coincide with the minima of the potential of the uplifting sector before the combination, $V_B(\phi)$, and in general it has to be checked case by case. In the special case of uplifting to Minkowski we just argued that the critical point is stable, or marginally stable along z^α directions, so it is evident that the problem of uplifting the SUSY vacua of the $z-$sector to Minkowski has now reduced to finding the stable Minkowski minima of the uplifting sector. The conditions for the existence of SUSY-breaking Minkowski vacua have been extensively analyzed by Gomez-Reino and Scrucca [13, 44, 45] as well as in [43].

Finally, in these Minkowski backgrounds, the gravitino mass after uplifting is related to the gravitino mass of the uplifting sector alone by

$$m_{3/2}^2 = e^{A(z_0)} m_{3/2,\phi}^2 \tag{3.16}$$

which is a special case of the more general relation

$$e^{G(z_0,\phi_0)} = e^{A(z_0)} e^{B(\phi_0)} \tag{3.17}$$

3.2 Supersymmetric critical points; BPS configurations

As we discussed in the previous subsection, when the z^α fields are stabilized at a SUSY critical point of the $z-$sector, $\partial_\alpha A|_{z=z_0} = 0$, there is no contribution from this sector to SUSY breaking in the full theory at tree level, i.e. the $F-$terms associated to these fields vanish. Therefore, for the complete theory to be at a SUSY critical point we just have to impose the additional condition that the F-terms for ϕ also vanish: $|F_\phi|^2 = e^G G^{i\bar{j}} G_i G_{\bar{j}}|_{\phi=\phi_0} = 0$, which is satisfied if and only if the ϕ^i fields are stabilized at SUSY critical point of the $\phi-$sector, $G_i|_{\phi=\phi_0} = B_i(\phi_0) = 0$. In other words, after fixing the z-fields at the SUSY critical point of the z-sector, the remaining effective theory for the ϕ fields gives the correct information about the critical points *of the full theory.*

This is closely related to the idea of integrating out heavy chiral multiplets *supersymmetrically*, which was first considered in [41] (see also [46]) and it is no accident that they found the same ansatz (1.3). Suppose that the z^α fields belong to the heavy chiral multiplets we want to integrate out. Then, when the energy scale under consideration is much lower than the masses of the $z-$sector we can fix these fields to their v.e.v.'s to a good approximation. If the $z-$sector is stabilized at a SUSY critical point, the effective low energy theory for the $\phi-$sector has unbroken $\mathcal{N} = 1$ local supersymmetry and is described

by the Kähler function:

$$G_{eff}(\phi) \equiv G|_{z=z_0} = A(z_0) + B(\phi), \tag{3.18}$$

which according to our previous argument would give the correct SUSY critical points in the full theory, since $\partial_i G_{eff} = \partial_i B$.

Other supersymmetry properties are also correctly inferred from the "effective" ϕ-theory, for instance BPS configurations of the ϕ-sector are also BPS in the coupled theory since the z fields do not contribute to SUSY breaking.

3.3 Shift symmetries

Whenever the Kähler function has a shift symmetry[4], as $G(z+\bar{z})$, or $G(\phi+\bar{\phi})$ there is a flat direction in the potential. For example if we assume

$$(\partial_z - \partial_{\bar{z}})G = 0 \quad \text{we have} \quad (\partial_z - \bar{\partial}_{\bar{z}})V = 0 \quad , \qquad V = V(z+\bar{z}). \tag{3.19}$$

The ansatz (3.4) ensures that the shift symmetries of A or B are also shift symmetries of full Kähler function G. Then if one of the two sectors has a flat direction in the potential which is related to a shift symmetry in its Kähler function, the same flat direction will survive in the full potential. This statement is still true for an arbitrary number of coupled sectors.

We can find an example of this situation in [33]. Here Hsu *et al.* give an effective SUGRA description of D3-D7 brane inflation in a type IIB string compactification. The D3-D7 configuration is BPS, resulting in a supersymmetric flat direction of the potential, which corresponds to the distance between the D7 and the D3 branes. Such a flat direction was implemented by introducing a Kähler function with a shift symmetry:

$$G = -3\ln(\rho+\bar{\rho}) - \frac{(S-\bar{S})^2}{2} + \ln|W_{\text{KKLT}}(\rho)|^2 \tag{3.20}$$

where ρ is the volume modulus, $W_{\text{KKLT}} = W_0 + Ae^{-a\rho}$ is the KKLT potential and S is a modulus describing the relative distance between a probe D7 brane and a heavy stack of D3 branes. The scalar potential derived from this Kähler function is independent of $\text{Re}(S)$ as a consequence of the shift symmetry in G.

In order to be able to use the shift symmetry as an inflationary trajectory, first we have to uplift it to de Sitter. In ref [33] the uplifting was achieved with D-terms. If we want to do the uplifting with F-terms while preserving the shift symmetry at tree level we can find two possibilities. We can add a new sector to the theory that is responsible for the uplifting and couple it using the ansatz (1.3), or we can add a term $\Delta G(S-\bar{S})$ to the

[4] Note that we are considering shift symmetries of the full Kähler function G, and not just of the Kähler potential K, so that the statement is consistent with Kähler invariance.

Kähler function such that, the $S-$sector is the one playing the role of the uplifting sector, which has been coupled using (1.3). We have studied the first possibility in a toy model with a single modulus in the $z-$sector. The results can be found in subsection 4.3.

Ref. [47] also describes a way to obtain exactly flat inflationary trajectories at tree level where the vacuum energy is also $F-$term dominated, but unlike in our case, the flat direction is not associated with a shift symmetry of the Kähler function.

4. Moduli stabilization in a toy model

In this section we will study in detail the stability properties of a simple model consisting of one modulus field z and one or more uplifting fields ϕ^i, with Kähler function

$$G = A(z, \bar{z}) + B(\phi^i, \bar{\phi}^{\bar{i}}) \ . \tag{4.1}$$

We start by calculating the mass spectrum of the modulus at the critical point before we couple the uplifting sector. We then calculate the stability after uplifting to dS or Minkowski space. The conclusions are summarized in figure 1.

4.1 Stability of the critical point before uplifting

In order to find the mass spectrum we expand the potential around the supersymmetric critical point z_0, $z = z_0 + \hat{z}$:

$$V = V(z_0) + \frac{1}{2}V_{zz}(z_0)\ \hat{z}^2 + \frac{1}{2}V_{\bar{z}\bar{z}}(z_0)\ \hat{\bar{z}}^2 + V_{z\bar{z}}(z_0)\ \hat{z}\hat{\bar{z}} + \ldots, \tag{4.2}$$

The diagonalization of the mass matrix gives us the spectrum of masses squared:

$$m^2_{\pm} = (V_{z\bar{z}}(z_0) \pm |V_{zz}(z_0)|)/A_{z\bar{z}}(z_0). \tag{4.3}$$

The condition for a local minimum is therefore $V_{z\bar{z}}(z_0) > |V_{zz}(z_0)| > 0$. Now we will write this condition in terms of the Kähler function A. Using 2.6, with G replaced by A, we find for the first derivative of the potential:

$$V_z = A_z V + e^A(A_z^{z\bar{z}} A_z A_{\bar{z}} + A^{z\bar{z}} A_{zz} A_{\bar{z}} + A^{z\bar{z}} A_z A_{z\bar{z}}) \tag{4.4}$$

Then the second derivatives evaluated at the critical point $z = z_0$ read:

$$V_{zz}(z_0) \ = \ -A_{zz}(z_0)e^{A(z_0)} \tag{4.5}$$

$$V_{z\bar{z}}(z_0) \ = \ e^{A(z_0)}[A^{z\bar{z}}|A_{zz}|^2 - 2A_{z\bar{z}}]_{z=z_0} \tag{4.6}$$

Here we used the assumption that we are expanding around a critical point and thus $A_z(z_0) = A_{\bar{z}}(z_0) = 0$. Defining

$$x \equiv \left|\frac{A_{zz}}{A_{z\bar{z}}}\right|_{z=z_0} , \tag{4.7}$$

we find that before uplifting

$$m^2_{\pm} = e^{A(z_0)}(|x|^2 - 2 \pm |x|) \tag{4.8}$$

which gives a characterization of the critical points in terms of $|x|$[5] :

$$\begin{aligned} |x| &> 2 && \text{local AdS minimum} \\ 1 < |x| &< 2 && \text{AdS saddle point} \\ |x| &< 1 && \text{local AdS maximum} \end{aligned} \tag{4.9}$$

Local maxima in AdS are not necessarily unstable [42] but such stability information is not necessary for the present calculation.

4.2 Stability after uplifting.

Take now the Kähler function to be of the form (4.1). Then, as it was discussed in section 3, after coupling the uplifting sector z_0 remains a critical point of the full potential. Moreover the mass matrix has a block diagonal form, $V_{zi}(z_0) = V_{z\bar{i}}(z_0) = 0$, and therefore the stability properties of the supersymmetric sector A can be studied by just considering the derivatives of the potential w.r.t. z and $\bar{z}$, and the resulting stability condition for the field z is again of the form (4.3). We just have to calculate V_{zz} and $V_{z\bar{z}}$. From (3.6) we obtain:

$$V_{zz}|_{z=z_0} = [e^B V_{A\ zz} + A_{zz} e^A V_B + 3A_{zz} e^{A+B}]_{z=z_0} \tag{4.10}$$

$$V_{z\bar{z}}|_{z=z_0} = [e^B V_{A\ z\bar{z}} + A_{z\bar{z}} e^A V_B + 3A_{z\bar{z}} e^{A+B}]_{z=z_0}. \tag{4.11}$$

We can recast these equations in a more compact form using (4.7), and using the abbreviation (3.12) $b \equiv B^{i\bar{j}} B_i B_{\bar{j}}$, which is only a function of the uplifting sector:

$$V_{zz}|_{z=z_0} = e^{A+B}|_{z=z_0}(b-1)\, x \qquad V_{z\bar{z}}|_{z=z_0} = e^{A+B}|_{z=z_0}(|x|^2 + b - 2) \tag{4.12}$$

Here, as in the previous section $|x| = |A_{zz}/A_{z\bar{z}}|_{z=z_0}$. Finally we can write the spectrum of masses squared around the critical point:

$$m_\pm^2 = e^{A+B}|_{z=z_0}\Big[(|x|^2 + b - 2) \pm |(b-1)x|\Big] = e^{A+B}|_{z=z_0}\left[\left(|x| \pm \frac{1}{2}(b-1)\right)^2 - \frac{1}{4}(b-3)^2\right] \tag{4.13}$$

To simplify the mass formula we assumed that $b > 1$. In the opposite case, $b < 1$, the masses m_+^2 and m_-^2 are exchanged. The stability condition for the field in the supersymmetric sector after uplifting reads:

$$m_\pm^2 = e^{A+B}|_{z=z_0}\left[\left(|x| \pm \frac{1}{2}(b-1)\right)^2 - \frac{1}{4}(b-3)^2\right] \geq 0 \tag{4.14}$$

The solutions to these inequalities in the case of uplifting to Minkowski or de Sitter, $b \geq 3$, are presented in fig.(1). We list here some interesting properties:

- Notice that the stability properties do not depend on the details of the uplifting sector, just on the amount of uplifting b. This actually fits in the intuition of weakly coupled systems.

[5] Note that the stability condition is invariant under $x \to e^{i\theta}\, x$, which is related to the U(1) symmetry $z \to e^{-i\frac{\theta}{2}}\, z$.

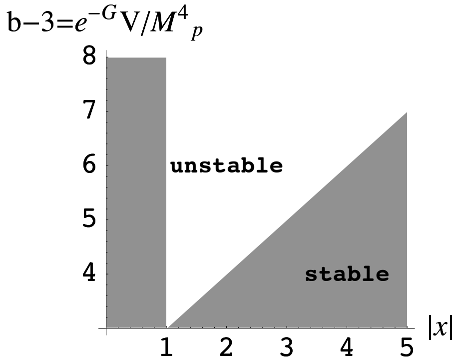

Figure 1: Stability of critical points after uplifting to Minkowski ($b = 3$) or de Sitter ($b > 3$) in the toy model described in the text. The shaded areas indicate stability along the moduli z directions. The vertical axis shows the quantity $b - 3 = e^{-G}V/M_p^4$ evaluated at the critical point, which represents the amount of uplifting. The horizontal axis shows the value of the quantity $|x| = |G_{zz}/G_{z\bar{z}}||_{z=z_0}$ at the critical point. Local minima before uplifting ($|x| > 2$) become unstable for sufficiently large upliftings. By contrast, local maxima before uplifting ($|x| < 1$) become more stable with increasing uplifting.

- All critical points z_0 that were local minima before the uplifting ($b = 0,\ |x| > 2$) remain stable for a certain amount of uplifting, and then all became unstable. As an example, the minimum found in the original KKLT paper [1] had $|x| \sim 25$.
- Critical points that are local maxima with $|x| < 1$ before uplifting, $b = 0$, become stable for $b = 3$, and remain stable for arbitrarily higher values of b. These points correspond to local minima of $e^{A/2}$.

In the case $x = 0$, which corresponds with having no $\hat{z}^2$ terms in $A(z, \bar{z})$, the two masses are equal, $m_{\pm}^2 = (b - 2)e^{A+B}|_{z_0}$, and both positive for $b > 2$. Points with $|x| = 1$ have one of the masses equal to zero for any uplifting.

Uplifting to *non-supersymmetric* Minkowski vacua has a special property. If $b = 3$ the mass squared

$$(m^2\pm)_{Minkowski} = m_{3/2}^2(|x| \pm 1)^2 \tag{4.15}$$

is positive definite for *any* choice of $A(z)$ (any value of $|x|$). Here we have used that in Minkowski vacuum the gravitino mass is given by $m_{3/2}^2 = e^{A+B}|_{z_0,\phi_0}$. This situation is close in spirit to the global susy case where critical points are always absolute minima. Here, Minkowski vacua are local minima of the supersymmetric sector whenever $|x| \neq 1$ or have a zero mode when $|x| = 1$. In the next section we give an explicit example of this latter case based on a shift symmetry.

4.3 A simple example: an uplifted flat direction in dS protected by shift symmetry

We will now consider the case where the Kähler function of the supersymmetric sector $A(z, \bar{z})$ has a shift symmetry, we will take it to be of the form $A = A(z + \bar{z})$. Given the

shift symmetry, ∂_z and $\partial_{\bar{z}}$ are interchangeable when they act on A or V so, in particular we have

$$A_z = A_{\bar{z}}, \quad A_{zz} = A_{z\bar{z}} \quad \text{and} \quad V_{zz} = V_{z\bar{z}}. \tag{4.16}$$

Suppose now that $A(z, \bar{z})$ has a SUSY critical point[6]. For this critical point we have $|x| = |A_{zz}/A_{z\bar{z}}| = 1$. Before uplifting there is one flat direction (zero mass) and one "tachyonic" direction with negative mass squared

$$m_-^2 = 0 \tag{4.17}$$

$$m_+^2 = 2V_{z\bar{z}}/A_{z\bar{z}}|_{z=z_0} = -2e^A|_{z=z_0} < 0. \tag{4.18}$$

The zero mode reflects the fact that the potential does not depend on Imz, and the Re z direction is always a local maximum since $A_{zz} = A_{z\bar{z}} > 0$, although not necessarily unstable since we are in AdS. After uplifting, the mass squared becomes positive while the flat direction remains

$$m_-^2 = 0 \tag{4.19}$$

$$m_+^2 = 2V_{z\bar{z}}/A_{z\bar{z}}|_{z=z_0} = e^{A+B}|_{z=z_0}(b-1), \tag{4.20}$$

so in this case it seems we can have positive mass squared whenever $b = B^{i\bar{j}}B_iB_{\bar{j}} \geq 1$, and in particular whenever $b \geq 3$. We note, however, that these results are only in our toy model but whether they generalize to the case with several moduli remains to be seen.

For $b > 3$ this simple model has a de Sitter, exactly flat z direction protected by the shift symmetry. Note that this "inflaton trench" is an F-term-uplifted AdS "ridge" (a line of local maxima), in contrast with the one proposed in [33], which was an AdS "trench" uplifted by D-terms. Its viability as an inflationary trajectory depends on whether the quantum corrections (from couplings to other fields) will tilt the flat direction to the required level. Alternatively, a soft breaking of the shift symmetry can be introduced. A graceful exit from inflation requires a more complicated scenario. But the point we want to emphasize is that there is no η problem.

5. Summary and discussion

Motivated by the KKLT uplifting problem, we have investigated a class of models where the Kähler potential and the superpotential are of the form:

$$K(z^\alpha, \bar{z}^{\bar{\alpha}}, \phi^i, \bar{\phi}^i) = K^{(1)}(z^\alpha, \bar{z}^{\bar{\alpha}}) + K^{(2)}(\phi^i, \bar{\phi}^i)$$
$$W(\phi^i, z^\alpha) = W^{(1)}(z^\alpha)W^{(2)}(\phi^i),$$

or, equivalently, where the full Kähler function is of the form

$$G = G^{(1)}(z^\alpha, \bar{z}^{\bar{\alpha}}) + G^{(2)}(\phi^i, \bar{\phi}^i) \ . \tag{5.1}$$

We have shown that these models have a number of interesting general properties:

[6]The KKLT superpotential for the volume modulus with $W_0 = 0$ is of this form, $W = Ae^{-az}$ which gives $G = -3\log(z+\bar{z}) - a(z+\bar{z}) + const.$ but its SUSY critical point is unphysical since it has negative $z + \bar{z}$ at z_0, $A_z(z_0) = 0$

- If $z^\alpha = z_0^\alpha$ is a SUSY critical point in the model defined by $G^{(1)}(z^\alpha, \bar{z}^{\bar{\alpha}})$, that is, if $(\partial G^{(1)}/\partial z^\alpha)|_{z^\alpha = z_0^\alpha} = 0$, this sector will not contribute to SUSY breaking in the full model: $F_z \propto (\partial G/\partial z^\alpha)|_{z^\alpha = z_0^\alpha} = 0$. Moreover $z^\alpha = z_0^\alpha$ is automatically a critical point of the combined (uplifted) potential in the z-direction.

- The stability of uplifted SUSY critical points of the $z-$sector, can be analyzed independently in the z^α and ϕ^i directions, since the crossed second derivatives of the combined potential vanish at this point: $\partial_\alpha \partial_i V|_{z=z_0} = \partial_{\bar{\alpha}} \partial_i V|_{z=z_0} = 0$.

- Local minima of the ϕ-potential when the uplifting sector is considered alone –the model defined by $G^{(2)}(\phi^i, \bar{\phi}^{\bar{i}})$–, *always* remain local minima in the ϕ-directions after the uplifting of the SUSY critical point of the $z-$sector, z_0^α. Moreover if the $\phi-$sector is stabilized at a Minkowski or de Sitter vacuum, z_0^α will be uplifted to Minkowski and de Sitter respectively.

- When a supersymmetric critical point is uplifted to Minkowski the critical point becomes automatically stable or flat along the z^α directions. A similar result was obtained in [43], where it was proven that all SUSY Minkowski critical points are stable. However our result describes the uplifting of SUSY critical points to *non supersymmetric* Minkowski vacua.

- When a supersymmetric critical point is uplifted to de Sitter, for sufficiently large cosmological constant the local minima of $G^{(1)}(z_\alpha, \bar{z}_\alpha)$ are always local minima of the combined potential (after uplifting) along the z^α-directions. Moreover, this local minimum becomes more stable with increasing value of the cosmological constant. Note that local minima of $G^{(1)}$ are always extrema of the moduli potential before uplifting, but not necessarily local minima.

- Shift symmetries of the individual sectors survive after the uplifting, becoming interesting candidates for inflationary trajectories.

We have studied in detail a toy model with a single field in the supersymmetric sector, where we have analyzed the stability of the $z-$sector "before", and "after" the uplifting. We have confirmed that uplifting to Minkowski space is special in that all SUSY critical points (irrespective of the choice of $G^{(1)}(z, \bar{z})$) become stable or neutrally stable. Indeed, after uplifting to Minkowski, the moduli masses are given by

$$m_\pm^2 = m_{3/2}^2 (|x| \pm 1)^2 \ , \quad \text{with} \quad |x| = |G^{(1)}_{zz}/G^{(1)}_{z\bar{z}}|_{z=z_0} = |G_{zz}/G_{z\bar{z}}|_{z=z_0}. \tag{5.2}$$

Note that if $|x| > 2$ the masses of the scalars in the supersymmetric sector are larger than the gravitino mass, for example, in the case of the KKLT model, $|x| \sim 25$, they are considerably larger. In general for $|x| < 1$, $m_\pm$ are of the order of the gravitino mass, except when the value of x is very close to 1, because in this case m_- becomes significantly lower than $m_{3/2}$. The case $|x| < 1$ is interesting because an uplifted critical point is stable for an arbitrary amount of uplifting. These critical points are precisely the minima of $G^{(1)}(z, \bar{z})$, which in this toy model correspond to local AdS maxima of the scalar potential before

uplifting.

Finally, we have shown that performing Kähler transformations before coupling two sectors gravitationally

$$K = K^{(1)} + K^{(2)} \qquad W = W^{(1)} + W^{(2)} \ , \tag{5.3}$$

may lead to direct couplings, and therefore the choice of Kähler gauge plays an important role in the applicability of this prescription.
An important consideration is whether string theory contains sectors that are coupled in the way described in this paper. We think it may be possible to find such couplings in certain $\mathcal{N} = 2$ compactifications where the presence of fluxes breaks supergravity down to $\mathcal{N} = 1$. $\mathcal{N} = 2$ supergravity requires that the kinetic terms of the scalars of vector and hypermultiplets appear totally decoupled from each other, and although the scalar manifold in general gets distorted during the SSB, there are known cases where this decoupling prevails [48, 49, 50, 51]. This leads to effective $\mathcal{N} = 1$ theories with a Kähler potential of the form (1.4). Moreover it is known that if the effective $\mathcal{N} = 1$ resulting from the SSB is consistent with a truncation of $\mathcal{N} = 2$ there are situations where the superpotential has a product structure

$$W = W^{(1)}(vector)\, W^{(2)}(hyper) \tag{5.4}$$

where $W^{(1)}$ and $W^{(2)}$ depend only on the scalars that came from the $\mathcal{N} = 2$ vector and hypermultiplets respectively [52, 53], which is precisely the kind of structure that we have studied in this paper.

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